

Emergent gauge theories and supersymmetry: a QED primer

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Abstract

We argue that a generic trigger for photon and other gauge fields to emerge as massless Nambu-Goldstone modes could be spontaneously broken supersymmetry rather than physically manifested Lorentz violation. We consider supersymmetric QED model extended by an arbitrary polynomial potential of vector superfield that induces the spontaneous SUSY violation in the visible sector. As a consequence, massless photon appears as a companion of massless photino being Goldstone fermion state in tree approximation. Remarkably, the photon masslessness appearing at tree level is further protected against radiative corrections due to the simultaneously generated special gauge invariance in the broken SUSY phase. Meanwhile, photino being mixed with another goldstino appearing from a spontaneous SUSY violation in the hidden sector largely turns into light pseudo-goldstino whose physics seems to be of special interest.

1 Introduction and overview

It is well known that spontaneous Lorentz invariance violation (SLIV) may lead to an emergence of massless Nambu-Goldstone modes [1] which are identified with photons and other gauge fields appearing in the Standard Model. This idea [2] supported by a close analogy with the dynamical origin of massless particle excitations for spontaneously broken internal symmetries has gained a new development [3, 4, 5, 6, 7] in recent years¹.

In this connection, one important thing to notice is that, in contrast to the spontaneous violation of internal symmetries, SLIV seems not to necessarily imply a physical breakdown of Lorentz invariance. Rather, when appearing in a gauge theory framework, this may eventually result in a noncovariant gauge choice in an otherwise gauge invariant and Lorentz invariant theory. In substance the SLIV ansatz, due to which the vector field $A_\mu(x)$ develops a vacuum expectation value (vev)

$$\langle A_\mu \rangle = n_\mu M \quad (1)$$

(where n_μ is a properly-oriented unit Lorentz vector, $n^2 = n_\mu n^\mu = \pm 1$, while M is the proposed SLIV scale), may itself be treated as a pure gauge transformation with a gauge function linear in coordinates, $\omega(x) = n_\mu x^\mu M$. From this viewpoint gauge invariance in QED leads to the conversion of SLIV into gauge degrees of freedom of the massless Goldstonic photon emerged. We will hereafter refer to it as an "inactive" SLIV, as opposed to an "active" SLIV leading to physical Lorentz violation (which may appear if gauge invariance is explicitly broken, see below).

A good example for such a kind of the inactive SLIV is provided by the nonlinearly realized Lorentz symmetry for underlying vector field $A_\mu(x)$ through the length-fixing constraint

$$A_\mu A^\mu = n^2 M^2. \quad (2)$$

This constraint in the gauge invariant QED framework was first studied by Nambu a long ago [11], and in more detail in recent years [12, 13, 14, 15, 16]. The constraint (2) is in fact very similar to the constraint appearing in the nonlinear σ -model for pions [17], $\sigma^2 + \pi^2 = f_\pi^2$, where f_π is the pion decay constant. Rather than impose by postulate, the constraint (2) may be implemented into the standard QED Lagrangian L_{QED} through an invariant Lagrange multiplier term

$$L = L_{QED} - \frac{\lambda}{2} (A_\mu A^\mu - n^2 M^2) \quad (3)$$

provided that initial values for all fields (and their momenta) involved are chosen so as to restrict the phase space to values with a vanishing multiplier function $\lambda(x)$, $\lambda = 0$ ².

¹Independently of the problem of the origin of local symmetries, Lorentz violation in itself has attracted considerable attention as an interesting phenomenological possibility which may be probed in direct Lorentz non-invariant extensions of quantum electrodynamics and the Standard Model [8, 9, 10].

²Actually, due to an automatic conservation of the matter current in QED an initial value $\lambda = 0$ will remain for all time. In a general case, when nonzero values of λ are also allowed, it appears problematic to have ghost-free theory theory with a positive Hamiltonian (for a detailed discussion see [18]). It is

One way or another, the constraint (2) means in essence that the vector field A_μ develops the vev (1) and Lorentz symmetry $SO(1,3)$ breaks down to $SO(3)$ or $SO(1,2)$ depending on whether the unit vector n_μ is time-like ($n^2 > 0$) or space-like ($n^2 < 0$). The point is, however, that, in sharp contrast to the nonlinear σ model for pions, the nonlinear QED theory, due to gauge invariance in the starting Lagrangian L_{QED} , ensures that all the physical Lorentz violating effects turn out to be non-observable. Actually, the nonlinear constraint (2) implemented as a supplementary condition appears in essence as a possible gauge choice for the vector field A_μ , while the S -matrix remains unaltered under such a gauge convention. Indeed, this nonlinear QED contains a plethora of Lorentz and CPT violating couplings when it is expressed in terms of the pure Goldstonic photon modes (a_μ) according to the constraint condition (2)

$$A_\mu = a_\mu + n_\mu(M^2 - n^2 a^2)^{\frac{1}{2}}, \quad n_\mu a_\mu = 0 \quad (a^2 \equiv a_\mu a^\mu). \quad (4)$$

However, the contributions of all these Lorentz violating couplings to physical processes completely cancel out among themselves. So, the inactive SLIV inspired by the length-fixing constraint (2) affects only the gauge of the vector potential A_μ while leaving physical Lorentz invariance intact, as was shown in tree [11] and one-loop [12] approximations. Later a similar result was also confirmed for spontaneously broken massive QED [13], non-Abelian theories [14] and tensor field gravity [16].

From this point of view, emergent gauge theories induced by the inactive SLIV mechanism are in fact indistinguishable from conventional gauge theories. Their Goldstonic nature could only be seen when taking the gauge condition of the length-fixing constraint type (2). As to their observational evidence, the only way for SLIV to cause physical Lorentz violation would appear if gauge invariance in the theory considered were really broken, rather than merely constrained by some gauge condition.

In this sense, a valuable alternative to the nonlinear QED model (3) might be a conventional QED type Lagrangian extended by an arbitrary vector field potential energy terms which explicitly break gauge invariance. For a minimal potential containing bilinear and quadrilinear vector field terms one comes to the Lagrangian

$$L = L_{QED} - \frac{\lambda}{4} (A_\mu A^\mu - n^2 M^2)^2 \quad (5)$$

where λ is now a coupling constant rather than the Lagrange multiplier field. This potential being sometimes referred to as the “bumblebee” model (see [7] and references therein) means in fact that the vector field A_μ develops a constant background value (1) causing again an appropriate (time-like or space-like) Lorentz violation at a scale M . However, in contrast to the nonlinear QED model with a directly imposed vector field constraint (3), the bumblebee model contains an extra degree of freedom which appears as a massive Higgs mode away from the potential minimum. This allows SLIV to get active. Indeed, due to the presence of this mode the model may lead to some physical Lorentz violation in

also worth noting that, though the Lagrange multiplier term formally breaks gauge invariance in the Lagrangian (3), this breaking, is in fact reduced to the nonlinear gauge choice (2) in the restricted phase space mentioned above.

terms of the properly deformed dispersion relations for photon and matter fields involved that appear from corresponding radiative corrections to their kinetic terms [5]. However, in either case, whether SLIV is active or inactive, it unavoidably leads to the generation of massless photons as vector Nambu-Goldstone bosons.

Nevertheless, it may turn out that SLIV is not the only reason why massless photons could dynamically appear, if spacetime symmetry is further enlarged. In this connection, a special interest may be related to supersymmetry. Actually, the situation is changed dramatically in the SUSY inspired models. It appears that, while the nonlinear QED model with its inactive SLIV successfully matches supersymmetry, the potential-extended QED models (including the bumblebee model discussed above) leading to physical Lorentz violation cannot be conceptually realized in the SUSY context³. This allows to think that physical Lorentz invariance is somewhat protected by SUSY and in this sense a generic trigger for massless photons to dynamically appear could be spontaneously broken supersymmetry itself rather than physically manifested SLIV. To see how this idea might work we consider supersymmetric QED model extended by an arbitrary polynomial potential of a general vector superfield that induces the spontaneous SUSY violation. As a consequence, massless photon emerges as a companion of massless photino being Goldstone fermion in the broken SUSY phase in the visible sector (Section 2). Remarkably, this masslessness appearing at tree level is further protected against radiative corrections by the simultaneously generated special gauge invariance in the Lagrangian at the SUSY breaking potential minimum (Section 3). Meanwhile, photino being mixed with another goldstino appearing from a spontaneous SUSY violation in the hidden sector largely turns into light pseudo-goldstino whose physics seems to be of special interest. Some overall conclusion is drawn in Section 4.

2 Extended supersymmetric QED

We now consider the supersymmetric QED extended by an arbitrary polynomial potential of a general vector superfield $V(x, \theta, \bar{\theta})$ whose a pure vector component A_μ is usually associated with a photon. The corresponding Lagrangian can be written in the SUSY invariant form as

$$\mathcal{L} = L_{SQED} + \sum_{n=1} b_n V^n|_D \quad (6)$$

where terms in this sum (b_n are some constants) for a conventional vector superfield parametrization⁴ are given by corresponding D -term expansions $V^n|_D$ into the compo-

³The point is that, in contrast to an ordinary vector field theory where all kinds of terms with any power of the vector field squared, $(A_\mu A^\mu)^n$ ($n = 1, 2, \dots$), can be included into the Lagrangian in Lorentz invariant way, SUSY theories only admit the bilinear term $A_\mu A^\mu$ in the vector field potential energy. This can readily be confirmed in SUSY QED extended by an arbitrary polynomial potential of a general vector superfield (Section 2). As a result, without a stabilizing high-linear (at least quadrilinear, as in (5)) vector field terms, the potential-based SLIV never occurs in SUSY theories.

⁴This parametrization is given by [19] $V(x, \theta, \bar{\theta}) = C(x) + i\theta\chi(x) - i\bar{\theta}\bar{\chi}(x) + \frac{i}{2}\theta\theta S(x) - \frac{i}{2}\bar{\theta}\bar{\theta}S^*(x) - \theta\sigma^\mu\bar{\theta}A_\mu(x) + i\theta\theta\bar{\theta}\bar{\lambda}'(x) - i\bar{\theta}\bar{\theta}\theta\lambda'(x) + \frac{1}{2}\theta\theta\bar{\theta}\bar{\theta}D'(x)$, where $S = M + iN$, $\lambda' = \lambda + \frac{i}{2}\sigma^\mu\partial_\mu\bar{\chi}$ and $D' = D + \frac{1}{2}\partial^2 C$.

nent fields. It can readily be checked that the first term in this expansion is the known Fayet-Iliopoulos D -term, while other terms only contain bilinear, trilinear and quadrilinear combination of the superfield components A_μ , S , λ and χ , respectively⁵. Actually, the higher-degree terms only appear for the scalar field component $C(x)$. Expressing them all in terms of the C field polynomial

$$P(C) = \sum_{n=1}^n \frac{b_n}{2} C^{n-1}(x) \quad (7)$$

and three of its derivatives with respect to the C field

$$P'_C \equiv \frac{\partial P}{\partial C}, \quad P''_C \equiv \frac{\partial^2 P}{\partial C^2}, \quad P'''_C \equiv \frac{\partial^3 P}{\partial C^3} \quad (8)$$

one has for the whole Lagrangian \mathcal{L}

$$\begin{aligned} \mathcal{L} = & -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} + i\lambda\sigma^\mu \partial_\mu \bar{\lambda} + \frac{1}{2} D^2 \\ & + \left(D + \frac{1}{2} \partial^2 C \right) P + \left(\frac{1}{2} S S^* - \chi\lambda' - \bar{\chi}\bar{\lambda}' - \frac{1}{2} A_\mu A^\mu \right) P'_C \\ & + \frac{1}{2} \left(\frac{i}{2} \bar{\chi}\chi S - \frac{i}{2} \chi\chi S^* - \chi\sigma^\mu \bar{\chi} A_\mu \right) P''_C + \frac{1}{8} (\chi\chi\bar{\chi}\bar{\chi}) P'''_C \end{aligned} \quad (9)$$

where we properly calculated D -terms for the corresponding powers of the superfield $V(x, \theta, \bar{\theta})$ in the polynomial in (6) and collected like field couplings. Also, for more clarity, we still omitted matter superfields in the model (see some discussion below). As one can see, extra degrees of freedom related to the C and χ component fields in a general vector superfield $V(x, \theta, \bar{\theta})$ appear through the potential terms in (9) rather than from the properly constructed supersymmetric field strengths, as is only appeared for the vector field A_μ and its gaugino companion λ .

Note that all terms in the sum in (6) except Fayet-Iliopoulos D -term explicitly break gauge invariance. However, as we will see in the next section, the special gauge invariance constrained by some gauge condition will be recovered in the Lagrangian in the broken SUSY phase. Furthermore, as is seen from (9), the vector field A_μ may only appear with bilinear mass term in the polynomially extended superfield Lagrangian (6). This means that, in contrast to the non-SUSY theory case (where, apart from the vector field mass term, some high-linear stabilizing terms necessarily appear in similar polynomially extended Lagrangian), Lorentz invariance is still preserved. Actually, only supersymmetry appears to be spontaneously broken, as mentioned above.

Indeed, varying the Lagrangian \mathcal{L} with respect to the D field we come to

$$D = -P(C) \quad (10)$$

⁵Without loss of generality, we may restrict ourselves to the third degree superfield polynomial in the Lagrangian \mathcal{L} (6) to eventually have a theory with dimensionless coupling constants for component fields. However, for completeness sake, it seems better to proceed with a general case.

that finally gives the following potential energy for the field system considered

$$U(C) = \frac{1}{2}[P(C)]^2 . \quad (11)$$

being solely determined by the polynomial of the scalar field component $C(x)$ of the superfield $V(x, \theta, \bar{\theta})$. The potential (11) may lead to the spontaneous SUSY breaking provided that the polynomial P (7) has no real roots, while its first derivative has,

$$P \neq 0 , \quad P'_C = 0 . \quad (12)$$

This requires $P(C)$ to be an even degree polynomial with properly chosen coefficients b_n in (7) that will force its derivative P'_C to have at least one root, $C = C_0$, in which the potential (11) is minimized. Therefore, supersymmetry is spontaneously broken and the C field acquires the vev

$$\langle C \rangle = C_0 , \quad P'_C(C_0) = 0 . \quad (13)$$

As an immediate consequence, that one can readily see from the Lagrangian \mathcal{L} (9), a massless photino λ being Goldstone fermion in the SUSY broken phase make all the other component fields in the superfield $V(x, \theta, \bar{\theta})$, including the photon, to also become massless. However, the question then arises whether this masslessness of photon will be stable against radiative corrections⁶, since gauge invariance is explicitly broken in the Lagrangian (9). We show in the next section that it could be the case if the vector superfield $V(x, \theta, \bar{\theta})$ would appear to be properly constrained.

Before proceeding further, note that we have not yet included matter superfields in the model. In their presence, the spontaneous SUSY breaking in the visible sector we have used above should be properly combined with a spontaneous SUSY violation in the hidden sector to evade the supertrace sum rule [19] for masses of basic fermions and their superpartners. Actually, this sum rule is acceptably relaxed when taking into account large radiative corrections to masses of supersymmetric particles, as typically appears in gauge-mediated SUSY breaking models [19]. This changes the simplified picture discussed above: the strictly massless fermion eigenstate appears to be some mix of the visible sector photino λ and the hidden sector goldstino rather than the pure photino state. In the supergravity context, one linear combination of them is eaten through the super-Higgs mechanism to form the longitudinal component of the gravitino, while their orthogonal combination, which may be referred to as a pseudo-goldstino, gets some mass⁷. One may generally expect that SUSY is much stronger broken in the hidden sector than in the visible one that means the pseudo-goldstino is largely given by the pure photino state λ . These states seem to be of special observational interest in the model that, apart from some indication of the QED emergence nature, may shed a light on the SUSY breaking physics.

⁶These corrections include as ordinary corrections appearing in a conventional supersymmetric QED, so corrections related to new degrees of freedom emerging in a general vector superfield $V(x, \theta, \bar{\theta})$ in terms of the revived component fields C and χ . All these corrections appear relativistically invariant since, as was mentioned above, physical Lorentz invariance is preserved in the basic Lagrangian \mathcal{L} (9).

⁷The possibility that the Standard Model visible sector might also break supersymmetry thus giving rise to similar pseudo-goldstino states was also considered, though in a different context, in [20, 21].

3 Constrained vector superfield

We have seen above that the vector field A_μ may only appear with bilinear mass terms in the polynomially extended Lagrangian (9). Hence it follows that potential-based models, particularly the bumblebee model mentioned above (5) with nontrivial vector field potential containing both a bilinear mass term and a quadrilinear stabilizing term, can in no way be realized in the SUSY context. Meanwhile, the nonlinear QED model, as will become clear below, successfully matches supersymmetry.

Let us notice first that instead of gauge symmetry broken in the extended QED Lagrangian (9) some special gauge invariance is recovered in (9) at the SUSY breaking minimum of the potential (11). We will expand the action around the vacuum (13) by writing

$$C(x) = C_0 + c(x) \quad (14)$$

that gives for the C field polynomial $P(C)$ (7) and its derivatives (8) to the lowest order in the Higgs-like field $c(x)$

$$\begin{aligned} P(C) &\simeq P(C_0) + \frac{1}{2}P_C''(C_0)c^2, \quad P_C'(C) \simeq P_C''(C_0)c, \\ P_C''(C) &\simeq P_C''(C_0) + P_C'''(C_0)c, \quad P_C'''(C) \simeq P_C'''(C_0) + P_C''''(C_0)c \end{aligned} \quad (15)$$

with $P_C'(C_0) = 0$ taken at the minimum point, as is determined in (12). Now, combining the equations of motion for $c(x)$ and for some other component field, say $S(x)$, both derived by varying the Lagrangian (9), one has

$$\frac{1}{2}SS^* - \chi\lambda' - \bar{\chi}\bar{\lambda}' - \frac{1}{2}A_\mu A^\mu = O(c, c\partial^2 c), \quad \chi\chi = O(c) \quad (16)$$

where we have used approximate equalities (15) with typically sizeable values of all $P(C_0)$, $P_C''(C_0)$, $P_C'''(C_0)$, $P_C''''(C_0)$ taken at the minimum point C_0 . So, at the SUSY breaking minimum ($c \rightarrow 0$) we come to the constraints which are put on the V superfield components

$$C = C_0, \quad \chi = 0, \quad A_\mu A^\mu = |S|^2 \quad (17)$$

that also determine the corresponding D -term (10), $D = -P(C_0)$, for the spontaneously broken supersymmetry.

Another, more exact, way to keep the whole theory at the SUSY breaking minimum, thus eliminating Higgs-like mode $c(x)$ forever, is to properly constrain the vector superfield $V(x, \theta, \bar{\theta})$ from the outset. This can be done, by analogy with constrained vector field in the nonlinear QED model (3), through the following SUSY invariant Lagrange multiplier term

$$\mathcal{L}' = \mathcal{L} + \frac{1}{2}\Lambda(V - C_0)^2|_D \quad (18)$$

where $\Lambda(x, \theta, \bar{\theta})$ is some auxiliary vector superfield, while C_0 is the constant background value of the C field for which potential $U(C)$ has the SUSY breaking minimum (13).

We further find for the Lagrange multiplier term in (18) that (denoting $\tilde{C} \equiv C - C_0$)

$$\begin{aligned} \Lambda V^2|_D = & C_\Lambda \left[\tilde{C} D' + \left(\frac{1}{2} S S^* - \chi \lambda' - \bar{\chi} \bar{\lambda}' - \frac{1}{2} A_\mu A^\mu \right) \right] \\ & + \chi_\Lambda \left[2\tilde{C} \lambda' + i(\chi S^* + i\sigma^\mu \bar{\chi} A_\mu) \right] + \bar{\chi}_\Lambda [2\tilde{C} \bar{\lambda}' - i(\bar{\chi} S - i\chi \sigma^\mu A_\mu)] \\ & + \frac{1}{2} S_\Lambda \left(\tilde{C} S^* + \frac{i}{2} \bar{\chi} \chi \right) + \frac{1}{2} S_\Lambda^* \left(\tilde{C} S - \frac{i}{2} \chi \chi \right) \\ & + 2A_\Lambda^\mu (\tilde{C} A_\mu - \chi \sigma_\mu \bar{\chi}) + 2\lambda'_\Lambda (\tilde{C} \chi) + 2\bar{\lambda}'_\Lambda (\tilde{C} \bar{\chi}) + \frac{1}{2} D'_\Lambda \tilde{C}^2 \end{aligned} \quad (19)$$

where

$$C_\Lambda, \chi_\Lambda, S_\Lambda, A_\Lambda^\mu, \lambda'_\Lambda = \lambda_\Lambda + \frac{i}{2} \sigma^\mu \partial_\mu \bar{\chi}_\Lambda, D'_\Lambda = D_\Lambda + \frac{1}{2} \partial^2 C_\Lambda \quad (20)$$

are the standard⁴ component fields of the Lagrange multiplier superfield $\Lambda(x, \theta, \bar{\theta})$. Now varying the whole Lagrangian (18) with respect to these nondynamical fields and properly combining their equations of motion

$$\frac{\partial \mathcal{L}'}{\partial (C_\Lambda, \chi_\Lambda, S_\Lambda, A_\Lambda^\mu, \lambda_\Lambda, D_\Lambda)} = 0 \quad (21)$$

we come again to the constraints (17) imposed on the vector superfield components. As before in non-SUSY case (3), to have ghost-free theory with a positive Hamiltonian the initial values for all component fields of superfield $V(x, \theta, \bar{\theta})$ are chosen so as to restrict phase space to values with the vanishing component fields (20) of the multiplier superfield $\Lambda(x, \theta, \bar{\theta})$.

Remarkably, one way or another, we have come for the constrained vector superfield $V(x, \theta, \bar{\theta})$ to almost the same physical states, photino and photon, as in supergauge multiplet of conventional supersymmetric QED [19]. Actually, photino field λ appears to be constraint-free, while the vector potential A_μ is only constrained by the condition relating it to the nondynamical S field. One can now readily confirm that, as a consequence of the spontaneous SUSY violation inducing the constraints (17), some special gauge invariance is in fact recovered in the Lagrangian (6). First of all, this violation provides the tree level photon masslessness, as is clearly seen in the Lagrangian (9) when the potential minimum condition $P'_C(C_0) = 0$ is applied. The rest of gauge noninvariance caused by a general superfield polynomial in (6) is simply reduced to the nonlinear gauge choice $A_\mu A^\mu = |S|^2$ in a virtually gauge invariant theory since extra degrees of freedom in terms of the C and χ component fields are also eliminated. Taking the S field to be some constant background field we come to the SLIV constraint (2) being in an ordinary nonlinear QED discussed above. As is seen from this gauge, one may only have a time-like SLIV in the SUSY framework but never a space-like one. There also may be a light-like SLIV, if the S field vanishes⁸. All of them are in fact inactive SLIV models in which physical Lorentz invariance is left intact. So, any possible choice for the S field only leads to the particular gauge

⁸Indeed, this case, first mentioned in [11], may also mean spontaneous Lorentz violation with a nonzero $\text{vev} < A_\mu > = (\tilde{M}, 0, 0, \tilde{M})$ and Goldstone modes $A_{1,2}$ and $(A_0 + A_3)/2 - \tilde{M}$. The "effective" Higgs mode $(A_0 - A_3)/2$ can be then expressed through Goldstone modes so that the light-like condition $A_\mu^2 = 0$ is satisfied.

choice for the vector field A_μ in an otherwise gauge invariant theory. Thus, the massless photon appearing first as a companion of massless photino (being Goldstone fermion in the visible broken SUSY phase) remains massless due to this recovering gauge invariance in the emergent SUSY QED. At the same time, the "built-in" nonlinear gauge condition (17) gives rise to treat the photon as vector Goldstone boson induced by the inactive SLIV.

4 Conclusion

We have attempted to find some supersymmetric analogue of the emergent QED inspired by spontaneous Lorentz invariance violation (SLIV). As discussed above, SLIV in a vector field theory framework may be inactive as in the nonlinear QED model (3), or active as in potential-extended QED models (5) leading to physical Lorentz violation. However, in either case SLIV unavoidably leads to the generation of massless Nambu-Goldstone modes which are identified with photons and other gauge fields appearing in the Standard Model. Nevertheless, it may turn out that SLIV is not the only reason why massless photons could dynamically appear, if spacetime symmetry is further enlarged. In this connection, a special interest is related to supersymmetry that we argued in this note by the example of supersymmetric QED⁹. Actually, while there are a few papers in the literature on Lorentz noninvariant extensions of supersymmetric models (for some interesting ideas, see [22] and references therein), an emergent gauge theory in a SUSY context has been considered for the first time.

In substance, the situation is changed dramatically in SUSY models: between two basic SLIV versions mentioned above, SUSY inevitably chooses the inactive SLIV case. Indeed, while the nonlinear QED model with its inactive SLIV successfully matches supersymmetry, the potential-extended QED models with an actual Lorentz violation can never occur in the SUSY context. This allows to think that a generic trigger for massless photons to appear could be spontaneously broken supersymmetry itself rather than physically manifested spontaneous Lorentz violation. For an explicit demonstration we considered supersymmetric QED model extended by an arbitrary polynomial potential of massive vector superfield that induces the spontaneous SUSY violation. As a consequence, massless photon emerges as a companion of massless photino being Goldstone fermion in the broken SUSY phase.

The photon masslessness appearing at tree level is further protected against radiative corrections by the simultaneously generated special gauge invariance. This invariance is only restricted by the nonlinear gauge condition put on vector field values, $A_\mu A^\mu = |S|^2$ with the nondynamical S field chosen as some arbitrary constant background field (including the vanishing one) in the theory. The point, however, is that this nonlinear gauge condition happens at the same time to be the SLIV type constraint which treats in turn the physical photon as the Lorentzian Goldstone mode. So, figuratively speaking, the photon passes through three evolution stages being initially the massive vector field component of a general vector superfield (9), then the three-level massless companion

⁹It is worth noting that all the basic arguments related to the present QED example can be extended to the Standard Model.

of the Goldstonic photino λ in the broken SUSY stage (13) and finally the generically massless state as the emergent Lorentzian mode in the inactive SLIV stage (17).

Meanwhile, photino being mixed with another goldstino appearing from a spontaneous SUSY violation in the hidden sector largely turns into light pseudo-goldstino whose physics seems to be of special interest. As argued in [21], if the SUSY visible sector possesses R -symmetry then the pseudo-goldstino mass is protected up to (R -violating) supergravity effects, and eventually the same region of parameter space simultaneously may solve both gravitino and pseudo-goldstino overproduction problems in the early universe. Apart from cosmological implications, many other sides of new physics related to pseudo-goldstinos appearing through the multiple SUSY breaking were also studied recently (see [20, 21, 23] and references therein). The point is, however, that there have been exclusively used non-vanishing F -terms as the only mechanism of visible SUSY breaking in models considered. In this connection, our pseudo-goldstonic photinos caused by non-vanishing D -terms in the visible SUSY sector may lead to somewhat different observational consequences. We are going to return to this interesting issue elsewhere.

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